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TeV Scale Left-Right Symmetry and Flavor Changing Neutral Higgs Effects

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In minimal left-right symmetric models, the mass of the neutral Higgs field mediating tree-level flavor changing effects (FCNH) is directly related to the parity breaking scale. Specifically, the lower bound on the Higgs mass coming from Higgs-induced tree-level effects, and exceeding about 15 TeV, would tend to imply a W_R mass bound much higher than that required by gauge exchange loop effects – the latter allowing W_R masses as low as 2.5 TeV. Since a W_R mass below 4 TeV is accessible at the LHC, it is important to investigate ways to decouple the FCNH effects from the W_R mass. In this Letter, we present a model where this happens, providing new motivation for LHC searches for W_R in the $1-4$ TeV mass range.

1. INTRODUCTION

Left-right symmetric models (LRSM), initially proposed to explain the origin of parity violation [1], have received further motivation from the fact that they provide a natural setting for understanding small neutrino masses via the seesaw mechanism [2]. The seesaw scale in these models happens to be the scale of parity breaking, thereby connecting the smallness of neutrino masses to the dominance of V−A interactions in low energy weak interactions. This connects two independently motivated new physics scales beyond the Standard Model (SM). Another compelling reason for considering this class of models is the Gell-Mann–Nishijima-like formula that relates electric charge, weak isospins and the baryon and lepton numbers of particles [3] as follows:

$$Q = T_{3L} + T_{3R} + (B-L)/2 \ , \qquad (1)$$

thereby justifying the hypercharge quantum numbers – that in the SM have just ad-hoc values – as remnants of the $SU(2)_R$ and $B-L$ gauge symmetries.

A question of great current interest is whether the new gauge bosons associated with the LRSM can be observed at the CERN LHC. Since the W_R has an unmistakable signature at the LHC (of two like sign dileptons and two jets, with no missing energy) [4], and a large enough production cross section if its mass is below 4 TeV [5], one needs to know any lower bounds on its mass from measured low energy observables. The relevant processes are the SM suppressed ones, in particular meson-antimeson mixing and CP violation in the kaon sector. Extensive analyses of these constraints have been carried out [6] over the years and the latest results can be stated as follows: the most stringent bounds come from the CP violating observables ϵ and ϵ' [6], and possibly the neutron electric dipole moment [7]; they however depend on how CP violation is incorporated into the model. For example, in the minimal versions of the model the left and right CKM angles are nearly the same, and the bounds from CP violating observables can even be absent altogether, if one defines parity as transforming $Q_L \to Q_L^c$ [8]. The next strongest bounds come, primarily, from the $K_L - K_S$ mass difference, and arise from new, gauge mediated contributions. Specifically, defining parity as $Q_L \to Q_R$, one finds a bound of 4 TeV [9], whereas with the $Q_L \to Q_L^c$ parity definition, one gets 2.5 TeV [8]. The weaker of the above bounds still allows the LHC to search for the W_R and the associated Z'.

There is however another source of flavor violation in left-right models [10]. Even in its minimal version, a LRSM contains two copies of the standard model Higgs doublet, embedded into the bi-doublet needed to generate fermion masses. This is therefore a version of two Higgs models which can be characterized, in the notation of [11], as having $Y_{2d} = Y_{1u}$ and $Y_{2u} = Y_{1d}$. As such, these models give rise to tree-level flavor changing neutral Higgs (FCNH) effects coming from the neutral member of the second SM Higgs doublet contained in the bi-doublet. We will call this the FCNH problem of LRSM. These effects turn out to put a lower limit on the second Higgs mass of around $10-15$ TeV. This bound would imply a corresponding bound on the parity breaking scale. In fact, the Higgs sector of minimal left-right models has been analyzed extensively [12] and it has been shown that, with the most general Higgs potential, the masses of the Higgs fields belonging to the second doublet owe their origin to the parity breaking scale M_{W_R}/g. Furthermore, if the Higgs self scalar couplings are kept in the perturbative range (say $\lambda_i \leq 1$), the FCNH constraints would raise the right handed W_R scale to the 10 TeV ballpark, pushing it out

of the LHC reach.

So the immediate question that arises is whether it is possible, and how naturally, to satisfy all the FCNH constraints while keeping the parity breaking scale accessible at the LHC. Some solutions to this problem have been suggested in the literature [13] by invoking supersymmetry and cancelling the above effects by supersymmetric box contributions or by using special CP violating solutions. Another route, where manifest left-right symmetry is forsaken, is to have the right handed quark-mixing angles different from the left handed ones [14]. The latter possibility (the so-called non-manifest LRSM), which has also been used to lower the bounds on the W_R mass itself [15], requires an extended Higgs sector.

In this Letter, we seek an alternative solution to the FCNH problem of LRSM within the non-supersymmetric framework via a minimal extension of the Higgs sector, without sacrificing the manifest nature of left and right handed quark mixings. Within our solution, the contributions to quark masses that break the relation $\mathrm{diag}(m_u, m_c, m_t) \propto \mathrm{diag}(m_d, m_s, m_b)$, as well as flavor mixing ($V^{\mathrm{CKM}} \neq \mathbb{1}$), both originate from an effective operator of dimension 5. It is straightforward to extend our model to the supersymmetric case.

The philosophy behind our approach is that the Higgs sector is independent of the gauge sector of generic models and their effects should in principle be controlled by separate physics. In particular, in the context of the left-right models, the Higgs sector should not be the determining factor as far as the scale of parity violation is concerned. Thus the FCNH problem of minimal LRSM may be a low energy manifestation of the fact that the model is incomplete and needs extension. Turning this question around, if indeed a low, TeV scale, W_R is discovered at the LHC, it would be an indication that the minimal LRSM needs extension in a way such that gauge and Higgs sectors are separated. We show in this Letter that it is indeed possible to separate the Higgs sector constraints from those of the gauge sector by introducing a discrete symmetry and a second bi-doublet with $B-L=2$ with mass M_ρ much higher than the left-right scale. The FCNH effects of the minimal left-right model are shown to be associated with this new mass scale, thus allowing the parity breaking scale to be determined solely by the W_R exchange effects at the one loop level. Although our goal is not to construct a grand unified theory, we note, incidentally, that the new Higgs bi-doublet in our model is part of the SO(10) multiplet **210** often used to break GUT symmetries. Below M_ρ and the scale of parity breaking, this model resembles a two Higgs doublet model such as the Glashow-Weinberg model [16] and therefore has natural flavor conservation. The suppression of FCNH effects can now be explained keeping all scalar couplings perturbative and without simultaneously "dragging" the W_R and Z' mass scale to the M_ρ bound.

This Letter is organized as follows: in sec. 2, we review the salient features of the minimal left-right symmetric model and the associated FCNH problem; in secs. 3 and 4, we present the new extended model and the mechanism for generating quark masses and mixings; in sec. 5, we discuss the associated FCNH effects and demonstrate the decoupling of the FCNH scale from the W_R scale. In sec. 6, we make some further remarks on the model.

2. THE LEFT-RIGHT SYMMETRIC MODEL: MOTIVATION AND MINIMAL REALIZATION

Left-right-symmetric models [1] are based on the group $G_{\mathrm{LR}} \equiv SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, with parity assumed to be a good symmetry of the Lagrangian. Let us first summarize their minimal realization. The SM matter content can be intuitively accommodated into fundamental representations of the G_{LR} group, namely

field	$\mathbf{G_{LR}}$ representation
$Q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}$	$(2,\ 1,\ 1/3)$
$Q_R = \begin{pmatrix} u_R \\ d_R \end{pmatrix}$	$(1,\ 2,\ 1/3)$
$L_L = \begin{pmatrix} \nu_L \\ \ell_L \end{pmatrix}$	$(2,\ 1,\ -1)$
$L_R = \begin{pmatrix} \nu_R \\ \ell_R \end{pmatrix}$	$(1,\ 2,\ -1)$

(2)

In particular, right handed neutrinos are introduced naturally by the requirement that right handed particles be also in doublets. These models therefore offer a very natural set-up for realizing the seesaw mechanism for small neutrino masses. The gauge currents associated with G_{LR} couple with strengths g_L, g_R and g', with $g_L = g_R$ to guarantee parity symmetry at higher energies.

The G_{LR} gauge group is broken spontaneously by the vev's of an appropriate Higgs sector. The Higgs sector most suitable for implementing the seesaw mechanism is given by a bi-doublet $\phi \sim (2,\ 2,\ 0)$ and two triplets $\Delta_L \sim (3,\ 1,\ 2)$ and $\Delta_R \sim (1,\ 3,\ 2)$, namely

$$\phi = \begin{pmatrix} \phi_1 & \phi_2 \end{pmatrix} = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} ,$$

$$\Delta_{L(R)} = \begin{pmatrix} \delta^+_{L(R)}/\sqrt{2} & \delta^{++}_{L(R)} \\ \delta^0_{L(R)} & -\delta^+_{L(R)}/\sqrt{2} \end{pmatrix} . \quad (3)$$

G_{LR} is broken first to $SU(2)_L \times U(1)_Y$ by the $\langle\Delta_R\rangle$ vev: we assume this scale to lie in the few TeV range. Electroweak symmetry breaking is then induced by the vev's of ϕ.

In left-right models, there are two different ways to define the parity operation: *(i)* $Q_L \leftrightarrow Q_R$, *(ii)* $Q_L \leftrightarrow Q_L^c$. The second definition arises naturally in the context of SO(10) grand unified models as well as supersymmetric left-right models. We will illustrate our new model using the first definition of parity although the results can be extended in a straightforward manner to the second case.

In the realization of parity invariance where $Q_L \leftrightarrow Q_R$, the four vev's associated with the ϕ and $\Delta_{L,R}$ fields can be complex: one can however perform [9] two field redefinitions, thereby setting two of the four phases to zero. Hence, in all generality, the Higgs vev's can be chosen as

$$\begin{aligned}\langle\phi\rangle &= \frac{1}{\sqrt{2}}\begin{pmatrix}\kappa & 0\\ 0 & \kappa' e^{i\alpha}\end{pmatrix}\,,\\ \langle\Delta_L\rangle &= \frac{1}{\sqrt{2}}\begin{pmatrix}0 & 0\\ v_L e^{i\theta_L} & 0\end{pmatrix}\,,\\ \langle\Delta_R\rangle &= \frac{1}{\sqrt{2}}\begin{pmatrix}0 & 0\\ v_R & 0\end{pmatrix}\,. \qquad (4)\end{aligned}$$

The quark fields and the Higgs bi-doublet in eqs. (2) and (3) give rise to the following Yukawa Lagrangian

$$\mathcal{L}_Y = \overline{\hat{Q}}_{Li}(h_{ij}\phi + \tilde{h}_{ij}\tilde{\phi})\hat{Q}_{Rj} + \text{h.c.} \qquad (5)$$

where $\tilde{\phi} = -i\tau_2\phi^* i\tau_2$, i, j are flavor indices and the hats indicate gauge eigenstates, to be distinguished from mass eigenstates. Parity symmetry, under which $Q_L \to Q_R$ and $\phi \to \phi^\dagger$, requires the Yukawa matrices h and $\tilde{h}$ to be Hermitian. Recalling that ϕ is a bi-doublet, it is clear that the terms on the r.h.s. of eq. (5) are all those present in the most general two Higgs doublet model (2HDM), except for the fact that the four Yukawa-coupling matrices allowed in the general case are here reduced to two, h and $\tilde{h}$, because of the left-right symmetry.

The Yukawa Lagrangian in eq. (5), after spontaneous symmetry breaking, gives rise to the quark mass matrices. Specifically, after performing the field redefinitions

$$\hat{U}_{L,R} = V^U_{L,R} U_{L,R}\,, \quad \hat{D}_{L,R} = V^D_{L,R} D_{L,R}\,, \qquad (6)$$

the diagonal mass matrices for up- and down-type quarks read

$$\begin{aligned}M_U &= \frac{1}{\sqrt{2}} V_L^{U\dagger}\left(\kappa h + e^{-i\alpha}\kappa'\tilde{h}\right)V_R^U\,,\\ M_D &= \frac{1}{\sqrt{2}} V_L^{D\dagger}\left(e^{i\alpha}\kappa' h + \kappa\tilde{h}\right)V_R^D\,. \qquad (7)\end{aligned}$$

The above mass matrices must reproduce, e.g., the relation $m_t \gg m_b$. Assuming no large cancellations between the two terms on either r.h.s. of eqs. (7), this implies the hierarchies $\kappa \gg \kappa'$ and $h \gg \tilde{h}$, namely that the first term in the expression for M_U will be dominant, $M_U \simeq \kappa h$, and one can always choose a basis where this matrix is diagonal. In this basis, the rotation on the d-quark fields that makes the M_D matrix diagonal would be

$$\hat{D}_{L,R} = V^{\rm CKM}_{L,R} D_{L,R}\,, \text{ with } V^{\rm CKM}_{L,R} = V^{U\dagger}_{L,R} V^D_{L,R}\,. (8)$$

One can easily convince oneself that, after the above rotations and after expressing the h and $\tilde{h}$ couplings in terms of the quark mass matrices, the U-U-Higgs (D-D-Higgs) couplings in eq. (5) that are proportional to the M_D (M_U) matrix, will be off-diagonal in flavor space because of the CKM misalignment, and give rise (among the other effects) to new tree-level contributions to flavor changing neutral current processes. In addition, the presence of the new phases implies new CP violating effects as well. In fact, a detailed analysis of these effects has been performed in ref. [9]. Meson mixings receive new loop contributions because of the W_R, entailing a lower bound on M_{W_R} around 4 TeV, or 2.5 TeV with an alternative parity definition. The already mentioned quark-quark-Higgs couplings do also contribute to these processes via tree-level diagrams mediated by the physical states H_1^0 and A_1^0. Assuming degeneracy in their masses, one obtains a lower bound of about 15 TeV. These bounds get in general even more severe once one takes into account also effects on CP violating observables, primarily ϵ_K. However, the CP violating effects are more model-dependent, and we omit them here for the sake of simplicity.

The above discussion is meant to highlight that, within the minimal LRSM realized as above, the lowest allowed parity breaking scale is typically determined by the FCNH mediated by the Higgs fields rather than by the new gauge boson exchanges, thereby pushing this scale out of the LHC range. In the next section we will show how the inclusion of an additional assumption is able to basically eliminate the above problems altogether, decoupling the parity breaking scale from the FCNH scale and keeping the parity breaking scale within reach of the LHC.

3. LRSM EXTENSION WITH A DISCRETE SYMMETRY

A. Definition

Looking at the Yukawa Lagrangian in eq. (5), it is evident that root of the severe FCNH problem inherent in the general LRSM is the mentioned presence of four Yukawa couplings, as they would appear in a general 2HDM [17], but locked in

two pairs, h and $\tilde{h}$, because of the LR symmetry. This interlocking makes it impossible to implement the minimal flavor violation hypothesis [18] in the LRSM. In order to suppress the FCNH effects, we therefore adopt the following procedure.

First, we note that the $\tilde{h}$ couplings would be forbidden in presence of a discrete symmetry, e.g. Z_4 defined as

$$\begin{aligned} &\phi \xrightarrow{Z_4} i\phi \ , \quad Q_R \xrightarrow{Z_4} -iQ_R \ , \\ &L_L \xrightarrow{Z_4} -iL_L \ , \quad \Delta_L \xrightarrow{Z_4} -\Delta_L \ , \end{aligned} \tag{9}$$

with all the other fields unchanged.[1]

In this case, only the coupling matrix $h \neq 0$, whereas $\tilde{h} = 0$, implying d-quark mass terms of the form $M_D = \kappa'/\kappa \, M_U$, which obviously cannot work for all the $m_{d,s,b}$ masses. Besides, the proportionality between M_D and M_U would not allow for the presence of a nontrivial CKM matrix for flavor or CP violation. We postulate that the observed u-quark mass patterns, along with the mismatch between the u- and the d-quark bases, are induced by additional effective operators of dimension ≥ 5. For this to occur, it is sufficient to assume the presence of a new kind of bi-doublet with non-zero $B-L$, denoted by $\rho \sim (2,\ 2,\ 2)$, namely

$$\rho = \begin{pmatrix} \rho_1^+ & \rho_2^{++} \\ \rho_1^0 & \rho_2^+ \end{pmatrix} \ , \quad \text{with } \rho \xrightarrow{Z_4} -i\rho \ . \tag{10}$$

As ρ is charged under $B-L$, it can couple to the $\Delta_{L,R}$ fields in a G_{LR}- (and Z_4-) invariant way, giving rise to dimension-5 operators. The Yukawa couplings for this model, invariant under Z_4, can be written as:

$$\begin{aligned} \mathcal{L}_Y \ &= \ \mathcal{L}_{Y,Q} \ + \ \mathcal{L}_{Y,L} \ , \\ \text{with} \qquad \mathcal{L}_{Y,Q} \ &= \ \overline{\hat{Q}}_L \, h\phi \, \hat{Q}_R \ + \ \frac{1}{M_\rho} \left(\overline{\hat{Q}}_L \, h_\rho \, \tilde{\rho}\Delta_R \, \hat{Q}_R \ + \ \overline{\hat{Q}}_L \, h'_\rho \, \Delta_L^\dagger \rho \, \hat{Q}_R \right) \ + \ L \leftrightarrow R \ , \end{aligned} \tag{11}$$

where all higher dimensional operators are suppressed by a new scale M_ρ. Taking into account the symmetry requirements (9), analogous terms can be written for the leptonic Yukawa Lagrangian, $\mathcal{L}_{Y,L}$, that is however not relevant to the rest of our discussion.

Note that under parity we assume that $\rho \to \tilde{\rho}^\dagger$ with all other definitions as usual (see e.g. [9]), so that all the Yukawa couplings can be made parity invariant. Below we analyze the effects on quark masses of adding these higher dimensional operators. In particular, since v_L is at most of the order of the left handed neutrino masses, we will drop the terms involving Δ_L in analyzing fermion masses.

B. Vacuum expectation value of the ρ field

A key ingredient of our discussion is that the ρ field acquires a vev and will contribute to quark masses. However, unlike the generic situation where the mass of the physical Higgs and its vev are of the same order, the ρ vev is induced by a tree-level tadpole, so that its mass is much larger, in the multi-TeV range, than its vev, instead of O(100 GeV). In fact, the discrete symmetry of our model allows a Higgs potential of the following form:

$$\begin{aligned} V(\phi, \Delta_R, \Delta_L, \rho) \ &= \ V_0(\phi, \Delta_R, \Delta_L) \\ &+ \ M_\rho^2 \mathrm{Tr}(\rho^\dagger \rho) \ + \ (M' \mathrm{Tr}(\phi^\dagger \tilde{\rho} \Delta_R) + \mathrm{h.c.}) \ , \end{aligned} \tag{12}$$

where V_0 is the general Higgs potential for ϕ, Δ_R and Δ_L fields discussed in ref. [12]. [2] The mass $M_\rho^2 > 0$ and minimization of the potential (12) with respect to the ρ vev gives

$$v_\rho = \frac{\kappa v_R}{\sqrt{2} M_\rho} \frac{M'}{M_\rho} \ , \tag{13}$$

where $\langle \rho^0 \rangle = v_\rho/\sqrt{2}$ and the other vevs are defined in eq. (4). As far as the value of v_ρ is concerned, eq. (13) shows that it depends on the dimensional coupling M' associated with the $\rho\phi\Delta_R$ term. In particular, depending on the ratio M'/M_ρ, the value of this vev can be either smaller or of order of the weak symmetry breaking scale. More details can be found in section 5 C.

[1] The possibility of the introduction of a horizontal symmetry, instead, has been discussed In the context of general (i.e. not LR symmetric) 2HDMs and in the LRSM in [19] and [20] respectively.

[2] See Appendix of Deshpande *et al.* in ref. [12]: in particular, $\mu_2, \lambda_4, \alpha_2 \to 0$ in view of our Z_4 symmetry, and the β_i terms are irrelevant to our discussion.

4. QUARK MIXINGS AND FCNH LAGRANGIAN

The Yukawa Lagrangian in eq. (11) gives rise to two contributions to quark masses. These contributions arise from the interactions of the neutral Higgs bosons, that get a vev. Three of them are relevant to our discussion: ϕ_1^0, ϕ_2^0 and ρ^0. Prior to symmetry breaking, their interactions can be written as:

$$\begin{aligned}\mathcal{L}_Y &= \overline{\hat{U}}_L\, h\phi_1^0\, \hat{U}_R \;+\; \overline{\hat{D}}_L\, h\phi_2^0\, \hat{D}_R \\ &- \frac{1}{M_\rho}\overline{\hat{U}}_L\, h_\rho\, \rho^{0*}\delta_R^0\, \hat{U}_R\ ,\end{aligned} \tag{14}$$

where, we recall, hats over the quark fields indicate flavor eigenstates. After symmetry breaking, the terms in eq. (14) give rise to the following mass terms

$$\begin{aligned} M_{\hat{U}} &= \frac{1}{\sqrt{2}}\left(h\kappa \;+\; \tilde{h}_\rho v_\rho\right)\ , \\ M_{\hat{D}} &= \frac{h\kappa'}{\sqrt{2}}\ ,\end{aligned} \tag{15}$$

where we introduced the coupling

$$\tilde{h}_\rho \equiv -\frac{h_\rho v_R}{\sqrt{2}M_\rho}\ . \tag{16}$$

Being in the flavor eigenbasis, the mass matrices $M_{\hat{U},\hat{D}}$ are still off-diagonal. We will denote the mass eigenbasis without the hat, and the diagonal mass matrices as $M_{U,D}$.

Without loss of generality, we can choose a basis where h is diagonal and hence so is the down quark mass matrix – so that $\hat{D} = D$. There is no further diagonalization of the down quark states. Therefore, the associated neutral Higgs boson ϕ_2^0 has only diagonal coupling to down quarks and as such does not lead to any flavor changing effects.

The CKM angles – and the corresponding flavor changing interactions – arise then from the couplings h and $\tilde{h}_\rho$ when the combination $h\kappa + \tilde{h}_\rho v_\rho$ is diagonalized to go to the mass eigenbasis for the up quarks. To study these interactions, we will henceforth neglect terms where the δ_R^0 field is dynamical, and only keep the Yukawa couplings proportional to its vev. The first step is to redefine the two neutral-Higgs components acting in the up-quark sector:

$$\begin{aligned} H_1^0 &= \frac{1}{\kappa_\rho}(\kappa\phi_1^0 + v_\rho\rho^{0*})\ , \\ H_2^0 &= \frac{1}{\kappa_\rho}(v_\rho\phi_1^0 - \kappa\rho^{0*})\ , \quad \text{with}\quad \kappa_\rho = \sqrt{\kappa^2 + v_\rho^2}\ ,\end{aligned} \tag{17}$$

such that H_1^0 and H_2^0 are orthogonal to each other, and $\langle H_2^0\rangle = 0$. The advantage of this basis is that, when we diagonalize the up-quark mass matrix, all the FCNH effects reside in the H_2^0 coupling, whereas the H_1^0 coupling becomes diagonal and does not contribute to FCNH effects.

In this basis, the neutral Higgs couplings for the up-quark mass eigenstates read

$$\mathcal{L}_{Y,U} = \frac{\sqrt{2}}{\kappa_\rho}\overline{U}_L\left[M_U\left(H_1^0 - \frac{\kappa}{v_\rho}H_2^0\right)\right]U_R \;+\; \frac{\sqrt{2}\kappa_\rho}{\kappa' v_\rho}\overline{U}_L\left(V_L^{\rm CKM}M_D V_R^{\rm CKM\dagger}H_2^0\right)U_R \;+\; {\rm h.c.}\ , \tag{18}$$

with $V_R^{\rm CKM} = V_L^{\rm CKM}$ because we are in a 'manifest' LR symmetric scenario, where CP is not violated spontaneously [9]. By substituting the $H_{1,2}^0$ vevs, one immediately sees that only the first term on the r.h.s. contributes to quark masses. Furthermore, only the second term on the r.h.s. is flavor off-diagonal. In the following section, we use this Lagrangian to discuss the FCNH effects and their decoupling from the W_R mass scale.

5. FCNH EFFECTS

A. Flavor changing Higgs admixture

From eq. (18) it is clear that FCNH effects decouple with the mass of H_2^0, which is an admixture of ϕ_1^0 and ρ^0. We should therefore first discuss the neutral-Higgs mass matrix in the rotated basis.

Defining $\Phi = \{\phi_1^0, \rho^{0*}\}$ as having square-mass matrix $\Phi^\dagger M_\Phi^2\Phi$, from the Higgs potential in eq. (12) one finds

$$M_\Phi^2 = \begin{pmatrix} 2\lambda_1\kappa^2 + \frac{v_R^2}{2}\frac{M'^2}{M_\rho^2} & -\frac{M'v_R}{\sqrt{2}} \\ -\frac{M'v_R}{\sqrt{2}} & M_\rho^2\end{pmatrix}\ . \tag{19}$$

While in general $\phi_{1,2}^0$ mix as well, in the limit

of $\kappa \gg \kappa'$ relevant here (see sec. 5 C), we can ignore those mixing effects and write the above matrix by itself. The 'flavor' basis in eq. (17), $H = \{H_1^0, H_2^0\}$, is obtained from the Φ basis through the rotation

$$R\Phi = H \ , \quad \text{with } R = \frac{1}{\kappa_\rho}\begin{pmatrix} \kappa & v_\rho \\ v_\rho & -\kappa \end{pmatrix} \ , \tag{20}$$

hence the H basis has square-mass matrix $M_H^2 = R M_\Phi^2 R^T$. In particular, the mass of the H_2^0 – the particle mediating FCNH effects – reads

$$M_{H_2^0}^2 \ = \ M_\rho^2 + \frac{v_R^2}{r_M^2}\left(1 + \lambda_1 \frac{\kappa^2}{M_\rho^2} + \frac{1}{4r_M^2}\frac{v_R^2}{M_\rho^2}\right) \tag{21}$$

where we have defined $r_M = M_\rho/M'$, and λ_1 is one of the parameters of the Higgs potential (see footnote 2).

B. Data: $D^0 - \overline{D}^0$ mixing

Up-quark FCNH effects can be bounded through $D^0 - \overline{D}^0$ mixing. The coupling relevant to this process is the one appearing in the U-U-H_2^0 interaction in eq. (18), namely

$$h_{UUH_2^0} \ = \ \frac{\sqrt{2}\kappa_\rho}{\kappa' v_\rho} V_L^{\rm CKM} M_D V_R^{\rm CKM\dagger} \ . \tag{22}$$

This coupling contributes to $D^0 - \overline{D}^0$ mixing via tree-level diagrams with exchange of H_2^0. A naive, order-of-magnitude, evaluation of these diagrams can be made by using the Goldstone theorem and the vacuum saturation approximation. The result is [3]

$$\Delta M_D \ \approx \ 2 \times \frac{(h_{UUH_2^0})_{12}^2}{M_{H_2^0}^2} \cdot \frac{1}{4}\frac{m_D^3 f_D^2}{(m_u + m_c)^2} \ \lesssim \ \Delta M_D^{\rm exp} \ , \tag{23}$$

where we take $\Delta M_D^{\rm exp} \ = \ x/\tau \ \simeq \ 1.6 \cdot 10^{-14}$ GeV, using $x \approx 0.97 \cdot 10^{-2}$ [22, 23] and $\tau = 410 \cdot 10^{-15}\ s$ [24]. Note that in eq. (23) we are assuming the experimental result for x to be saturated by new-physics contributions. This approach is justified, given the very poor knowledge of the SM contribution to ΔM_D [25]. Our results would anyway barely change if we assumed, e.g., $\Delta M_D \lesssim 0.5 \cdot \Delta M_D^{\rm exp}$.

[3] The factor of 1/4 in eq. (23) comes simply from the normalization of the helicity projectors as $(1 \pm \gamma_5)/2$ in the LR×RL operator, which in turn enters twice. (Note instead that the contributions from the other helicity combinations cancel [21].) Needless to say, this calculation omits a number of effects that in general are important, in particular the RGE enhancement of pseudoscalar operators.

Plugging eq. (22) into eq. (23), and recalling that

$$\begin{aligned} (V_L^{\rm CKM} M_D V_R^{\rm CKM\dagger})_{12} &\simeq m_s \lambda = 0.022 \text{ GeV} \ , \\ \frac{1}{2}\frac{m_D^3 f_D^2}{(m_u + m_c)^2} &\simeq 0.1 \text{ GeV}^3 \ , \end{aligned} \tag{24}$$

the naive bound in eq. (23) can be rewritten as

$$\begin{aligned} v_R \ &\gtrsim \ \frac{2 m_s \lambda}{B_U}\frac{\kappa_\rho}{\kappa\kappa'}\frac{M_\rho}{M_{H_2^0}}\frac{M_\rho}{M'} \ , \\ \text{with } B_U &= \sqrt{\frac{1.6 \cdot 10^{-14} \text{ GeV}}{0.1 \text{ GeV}^3}} \ . \end{aligned} \tag{25}$$

In the next section, we will discuss a refined version of this bound in the context of our numerical analysis. However, the bound in eq. (25) turns out to be very accurate. To get a numerical idea of this bound, taking $\kappa/\kappa' \approx 35$ and $\kappa_\rho \approx \kappa$, one would obtain

$$v_R \ \gtrsim \ (15 \text{ TeV}) \frac{M_\rho}{M_{H_2^0}}\frac{M_\rho}{M'} \ . \tag{26}$$

The crucial point here is that the mass ratios on the r.h.s. of eq. (25) can very easily provide a suppression factor of O(10) or even larger, so that v_R, and hence $M_{W_R} = g_R v_R$, is in the ballpark of 1 TeV.

C. Numerical Analysis

In order to test quantitatively the mechanism described above, we will now carry out a numerical exploration of the allowed parameter space of the model. To this end, the discussion in the previous sections allows to identify the following requirements:

1. The possibility we are mostly interested in, is that the scale of LR-symmetry breaking, v_R, be within LHC reach. Therefore we will enforce [5]

$$M_{W_R} = g_R v_R \lesssim 4 \text{ TeV} \ , \tag{27}$$

with $g_R = g_L$ identified with the SM SU$(2)_L$ coupling, implying $g_R \simeq 0.65$.

2. Barring accidental cancellations, which are unlikely in view of the hierarchical structure of the CKM matrix, one expects $(M_{\hat{U}})_{33} \,/\, (M_{\hat{D}})_{33}$ (see eq. (15)) to be of order $m_t/m_b \approx 40$. Therefore, we shall require

$$\Delta_{tb} \ \equiv \ \frac{\kappa}{\kappa'}\left(1 \pm \left|\frac{(h_\rho)_{33}}{(h)_{33}}\right| \left(\frac{v_R}{\sqrt{2}M_\rho}\right)^2 \frac{M'}{M_\rho}\right) \ = \ {\rm O}(m_t/m_b) \ . \tag{28}$$

Taking into account the many uncertainties, we consider the range $\Delta_{tb} \in [27, 60]$ a reasonable choice.

3. Our discussion is based on the quark–quark–neutral-Higgs Lagrangian in eq. (14). In order to justify the negligibility of contributions from operators of dimension higher than 5 allowed by the symmetry (9), we should require that the ratio between the dimension-5 and the dimension-4 contributions to $M_{\hat{U}}$ (see eq. (15)) be small enough. Specifically, we will take

$$\Delta_5 \equiv \frac{(\tilde{h}_\rho)_{33} v_\rho}{(h)_{33}\kappa} = \frac{(h_\rho)_{33}}{(h)_{33}} \frac{M'}{M_\rho} \left(\frac{v_R}{\sqrt{2} M_\rho} \right)^2 \lesssim 0.5 \; . \quad (29)$$

Values of Δ_5 even sensibly below this bound will turn out to be very easy to achieve, as shown in the plots to follow.

4. Finally, we shall enforce the FCNH bound from $D^0 - \overline{D}^0$ mixing. We have implemented a detailed FCNH calculation of ΔM_D, including running effects from $M_{H_2^0}$ to m_c, etc., and using inputs from refs. [26]. The final result reads [4]

$$\Delta m_D = \frac{(h_{UUH_2^0})^2_{12}}{M^2_{H_2^0}} \times 0.502 \text{ GeV}^3 \; , \quad (30)$$

displaying a substantial enhancement with respect to the naive bound in eq. (23). As anticipated in footnote 3, this enhancement is expected from RGE running effects between the $M_{H_2^0}$ and m_c scales. In fact, we have checked that, identifying the Wilson coefficients at the matching scale with those at the charm scale, we get back basically the same result as eq. (23). We will use eq. (30) as our FCNH constraint.

It is clear that the main model parameters are

$$\mathcal{P} \equiv \{v_R, M_\rho, M', \kappa\} \; , \quad (31)$$

the other vev's or vev combinations v_ρ, κ_ρ and κ' being fixed once the above parameters are. Note, in particular, that the identification of W_L with the SM W-boson allows to univocally set the κ' value as follows

$$M_W^{\text{SM}} = \frac{gv}{2} = M_{W_L} = g_L \frac{\sqrt{\kappa^2 + \kappa'^2}}{2} \; , \quad (32)$$

where $v \simeq 250$ GeV. Note as well that the constraints (28) and (29) together imply κ very close to v and κ' much smaller, of order κ/Δ_{tb}.

The mass scale of the flavor changing neutral-Higgs admixture is in principle a further free parameter, given that the coupling λ_1 (see eq. (21)) is free. However, we shall take $\lambda_1 \in [0.1, 20]$. With this requirement, the $M_{H_2^0}$ mass is again fixed once the set of parameters $\mathcal{P}$ is.

Finally, the constraints in eqs. (28) and (29) depend on the couplings ratio $(h_\rho)_{33}/(h)_{33}$. In order to absorb any suppression/enhancement effect into the dimensionful parameters of the model, we shall demand this couplings ratio to be broadly of order one. Conservatively we will require $(h_\rho)_{33}/(h)_{33} \in [0.1, 10]$.

The results of our numerical analysis are shown in fig. 1. The six panels display the M_{W_R} scale as a function of all the relevant parameters discussed above, namely M_ρ, M'/M_ρ, λ_1, $M_{H_2^0}$, v_ρ and the quantity Δ_5 defined in eq. (29).

Fig. 1 shows in the first place that M_{W_R} values in the range 1 to 4 TeV – well within the LHC reach [5] – are very natural to obtain. This requires M_ρ in the O(10 TeV) range, and values for M' and for $M_{H_2^0}$ higher than about $2M_\rho$, but not hierarchically higher. The fact that no fine tuning in these masses is needed is confirmed by the values for λ_1 – which, according to eq. (21), can be traded for $M_{H_2^0}$. The allowed λ_1 values in fig. 1 are quite uniformly distributed in their whole allowed range, for M_{W_R} in basically the entire interval $[1, 4]$ TeV.

In addition, the M_{W_R} vs. v_ρ panel shows (as a consistency check not imposed in the scan) that $v_\rho < v_R$ is always fulfilled, justifying in turn our approximation of neglecting FCNH effects due to the propagation of δ_R^0 with respect to those due to ρ^0.

Finally, as a further consistency check, the M_{W_R} vs. Δ_5 panel shows that the ratio Δ_5 (see definition in eq. (29)) between dimension-5 and dimension-4 contributions to the top mass is very naturally well below 1.

6. DISCUSSION

Here we collect some further remarks on the model.

(i) In the leptonic sector, not discussed in this Letter, the renormalizable ρ couplings must

[4] A straightforward way to obtain this result consists in implementing the (very convenient) formulae reported in Golowich *et al.*, ref. [26], in particular their eqs. (14), (82) and (83), with the strong coupling α_s evaluated at the different scales using `RunDec` [27]. Concerning α_s at the heavy-Higgs threshold, we chose $\alpha_s(10 \text{ TeV})$.

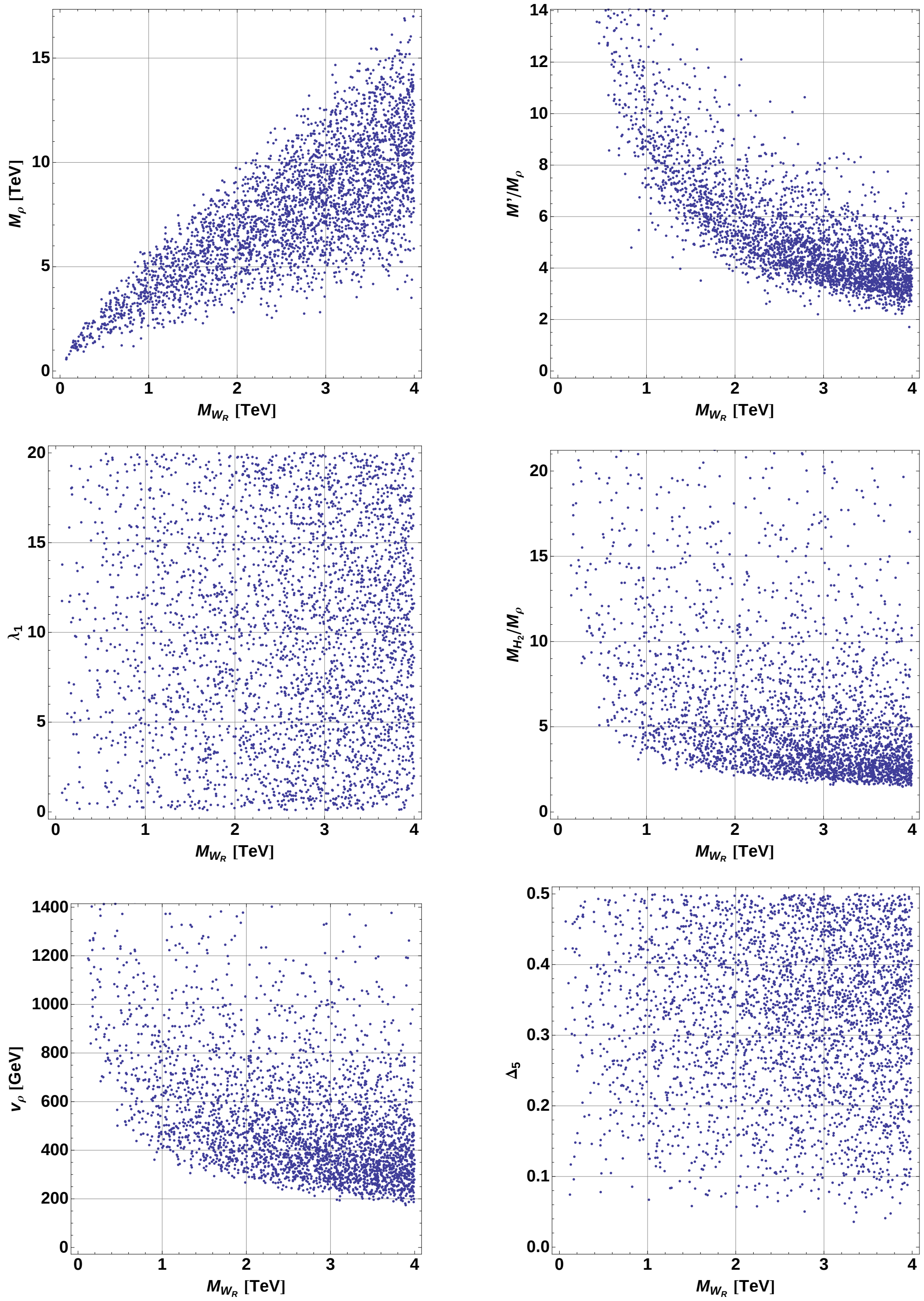

FIG. 1: Allowed parameter space for our model. See text, sec. 5 C, for comments on these plots.

be chosen fairly close to diagonal form, because all the neutrino mixings must arise from the right handed neutrino mass matrix to be consistent with lepton flavor violating branching ratios such as $\mu \to 3e$ etc.

(ii) Our idea can be extended to the supersymmetric left-right models and is particularly appealing in this case. The point is that, due to holomorphicity of the superpotential, in SUSYLR models one needs automatically two bi-doublets (in place of ϕ alone) to generate nontrivial CKM angles. Using our idea, one can now replace the second $B-L=0$ bi-doublet by a $B-L=2$ bi-doublet and use a higher dimensional operator of the

form $\frac{h_\rho}{M} Q^T \rho \Delta^c Q$ to generate CKM angles as well as to solve the FCNH problem while keeping the W_R scale in the LHC accessible range. No cancellations between Higgs exchange contribution and squark box graphs need be invoked [28].

(iii) The new ρ particles of the model are much heavier than the parity breaking scale and are therefore beyond the reach of the LHC.

(iv) The M' scale in the potential in eq. (12) could arise from the vev of a field σ reflecting higher scale physics, all our results remaining unchanged.

7. CONCLUSION

In summary, we have presented an extension of the left-right symmetric model which satisfies the bounds on Higgs masses arising from flavor changing effects, without at the same time dragging the parity breaking scale up with it. Within our model, the contributions to quark masses that break the relation diag$(m_u, m_c, m_t) \propto$ diag(m_d, m_s, m_b), as well as flavor mixing ($V^{\mathrm{CKM}} \neq \mathbb{1}$), are both the effect of a dimension 5 operator. This keeps the W_R and Z' within the reach of LHC and makes it possible, as the LHC collects data, to explore both the origin of parity violation and of neutrino masses – a common origin, in the context of our class of models. The added new particles are, on the other hand, beyond the LHC reach.

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